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On the mass correction of heavy meson effective theory *

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ABSTRACT

We derive the mass correction lagrangian of the heavy meson effective theory by using the projection operator method. The next leading order of the mass correction and the first order of the chiral expansion are given explicitely.We also consider the mass correction weak lagrangian of the heavy mesons. Finally, we give the $D^{*+} \to D^o \pi^+$ decay amplitude as an application.

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The heavy quark physics have gotten much progress in recent years. One reason comes from the discovery of the heavy quark symmetry of **QCD** by N. Isgur and M. B. Wise in **1989**[1]. This new symmetry appears in the limit of the infinite mass of the heavy quark , and can be used to relate different matrix elements of the heavy quark weak current. In **1990** , H. Georgi developed a heavy quark effective theory(**HQET**) to describe the above symmetry[2] for one heavy quark interacting with soft gluon by assuming the velocity superselection rule.

In **1991**, many authors continued the program of **HQET** to construct the heavy hadron effective theory (**HHET**) to combine heavy quark symmetry and chiral symmetry together[3]. As a member of the **HHET**, the heavy meson effective theory (**HMET**) considers one heavy meson interacting with the soft pseudoscalar pions[3].

Since the leading term in the lagrangian of **HMET** only holds in the infinite limit of the heavy quark mass , the mass correction effects need to be taken into account. The $1/m_Q$ correction has been studied by H. Y. Cheng, et al with interpolating field method and by N. Kitazawa and T. Kurimoto with velocity reparametrization invariant method[4]. Here we apply the projection operator method to derive the mass correction terms for **HMET** up to $1/m_Q^2$ order.

In this paper we construct the heavy meson effective lagrangian up to the $O(1/m_Q^2)$ corrections in $1/m_Q$ expansion and $O(p^2)$ in the chiral expansion. The lagrangian contains nineteen parameters(with six new couplings, three new mass splitting parameters between heavy pseudoscalar and vector

mesons)necessary to be determined by experiments.The light vector mesons are introduced via the hidden local symmetry method. The decay width of the process $D^{*+} \to D^o \pi^+$ is calculated with the $1/m_Q^2$ correction. Our notations are refered to [4].

We denote the compact heavy meson field with large but finite mass by $M_v(x)$, which is defined as

$$
\begin{aligned}
M_v(x) &= e^{i m_M v\cdot x} M_Q(x) \\
&= e^{i m_M v\cdot x} \left(\frac{i \not{\partial} + m_M}{2m_Q}\right) \left(-P_Q \gamma_5 + P_Q^{*\mu}\gamma_\mu\right)(x) \\
&= \left(\frac{i \not{\partial} + m_M + m_M \not{v}}{2m_Q}\right) \left(-\widehat{P}_v \gamma_5 + \widehat{P}_v^{*\mu}\gamma_\mu\right)(x) \\
&= \left(\frac{1+\not{v}}{2} + \frac{i \not{\partial}}{2m_Q} + \frac{\Lambda}{m_Q}\frac{1+\not{v}}{2}\right) \left(-\widehat{P}_v \gamma_5 + \widehat{P}_v^{*\mu}\gamma_\mu\right)(x),
\end{aligned}
\tag{1}
$$

where we have defined $m_M = m_Q + \Lambda$ and $M_Q(x)$ is the compact heavy meson field defined in [3]. The well-known field $H_v(x)$ is defined as the infinite mass limit of $M_v(x)$

$$
H_v(x) = \lim_{m_Q \to \infty} M_v(x) \tag{2}
$$

and the relevant ingredients of $H_v(x)$ are also taken as the same limit of those of $M_v(x)$

$$
\begin{aligned}
P_v(x) &= \lim_{m_Q \to \infty} \widehat{P}_v(x) \\
P_v^{*\mu}(x) &= \lim_{m_Q \to \infty} \widehat{P}_v^{*\mu}(x).
\end{aligned}
\tag{3}
$$

We take the pseudoscalar effective field $P_v(x)$ as example of application of the projection operator method to derive its mass correction form. Transforming (1) to its momentum space form and taking the pseudoscalar part , we then have

$$\begin{aligned}\left(\frac{1+\not v}{2}+\frac{\not k}{2m_Q}+\frac{\Lambda}{m_Q}\frac{1+\not v}{2}\right)\widehat{P}_v &= \left(\frac{1+\not v}{2}+\frac{\not k}{2m_Q}\right)\widehat{P}_v+\left(\frac{\Lambda}{m_Q}\frac{1+\not v}{2}\right)\widehat{P}_v \\ &\equiv \widehat{P}_{v,Q}+\widehat{P}_{v,\bar{q}}, \end{aligned} \tag{4}$$

where the $\widehat{P}_{v,Q}$ is the projected heavy degree of freedom and $\widehat{P}_{v,\bar{q}}$ is the projected light degree of freedom. Such a definition is similar to the Bargman-Wigner wave function in the Bathe-Salpeter approach found by F.Hussain et al in[6] but here they are the field operators.The residual momentum k comes from the momentum fluctuation of the heavy degree of freedom and thus is consistent with that definition of **HQET**. To arrive at the mass expansion formalism, we need to further separate $\widehat{P}_{v,Q}$ into $(1+\not v)/2$ and $(1-\not v)/2$ parts as

$$\widehat{\beta}_{v,Q}\equiv\left(\frac{1+\not v}{2}\right)\widehat{P}_{v,Q}\,;\,\widehat{\chi}_{v,Q}\equiv\left(\frac{1-\not v}{2}\right)\widehat{P}_{v,Q}. \tag{5}$$

We make an assumption that the heavy degree of freedom is on shell. Then the projection operator $((1+\not v)/2+\not k/2m_Q)$ becomes the energy projection operator(or, the spin projector) of the heavy quark. There will be the following constrain on the residual momentum k

$$v\cdot k=-\frac{k^2}{2m_Q}.$$

Applying $((1+\not v)/2+\not k/2m_Q)$ on $\widehat{P}_{v,Q}$, then we have

$$\frac{1-\not{v}}{2}\widehat{P}_{v,Q}=\frac{\not{k}}{2m_Q}\widehat{P}_{v,Q}, \tag{6}$$

and

$$\begin{aligned}\widehat{\chi}_{v,Q} &= \frac{\not{k}}{2m_Q}\left(\widehat{\beta}_{v,Q}+\widehat{\chi}_{v,Q}\right)\\ &= \frac{\not{k}}{2m_Q-\not{k}}\widehat{\beta}_{v,Q},\end{aligned} \tag{7}$$

These then give

$$\widehat{P}_{v,Q}=\frac{2m_Q}{2m_Q-\not{k}}\widehat{\beta}_{v,Q} \tag{8}$$

Therefore we only need to determine $\widehat{\beta}_{v,Q}$. By the follwing ansatz

$$\begin{aligned}\widehat{\beta}_{v,Q} &= (1+\widehat{\omega}_{v,Q})\,P_{v,h}\\ \overline{\widehat{\omega}}_{v,Q} &= \gamma^o\widehat{\omega}^{\dagger}_{v,Q}\gamma^o=\widehat{\omega}_{v,Q},\end{aligned} \tag{9}$$

and noting that

$$\Lambda^+_Q=\sum_r\widehat{P}_{r,v,Q}\overline{\widehat{P}}_{r,v,Q}\ ,\ \Lambda^+_h=\sum_r P_{r,v,h}\overline{P}_{r,v,h}, \tag{10}$$

with $\Lambda^+_Q=((1+\not{v})/2+\not{k}/2m_Q)$, $\Lambda^+_h=(1+\not{v})/2$, $P_{r,v,h}=([(1+\not{v})/2]P_v)_r$, we obtain

$$(1+\widehat{\omega}_{v,Q})\,\Lambda^+_h\,(1+\widehat{\omega}_{v,Q})=\left(1-\frac{\not{k}}{2m_Q}\right)\Lambda^+_Q\left(1-\frac{\not{k}}{2m_Q}\right). \tag{11}$$

(11) implies the following quadratical equation

$$\widehat{\omega}_{v,Q}^2 + 2\widehat{\omega}_{v,Q} - \widehat{T}_{v,Q} = 0, \tag{12}$$

where $\widehat{T}_{v,Q}$ is given by

$$\begin{aligned} \widehat{T}_{v,Q} &= \left(1 - \frac{\not{k}}{2m_Q}\right) \Lambda_Q^+ \left(1 - \frac{\not{k}}{2m_Q}\right) - \Lambda_h^+ \\ &= -\left(\frac{\not{k}}{2m_Q}\right)^2 \Lambda_h^+, \end{aligned} \tag{13}$$

and is of $O(m_Q^{-2})$. Since $\widehat{\omega}_{v,Q} \ll 1$, the relevant solution is

$$\begin{aligned} \widehat{\omega}_{v,Q} &= -1 + \left(1 + \widehat{T}_{v,Q}\right)^{1/2} \\ &= \frac{1}{2}\widehat{T}_{v,Q} - \frac{1}{8}\widehat{T}_{v,Q}^2 + \frac{1}{16}\widehat{T}3_{v,Q} - \frac{5}{128}\widehat{T}_{v,Q}^4 + \ldots. \end{aligned} \tag{14}$$

We note that the $\widehat{\omega}_{v,Q}$ in (14) agrees to the ansatz (9). By using

$$\left(\widehat{T}_{v,Q}\right)^n P_{v,h} = (-)^n \left(\frac{\not{k}}{2m_Q}\right)^{2n} P_{v,h}, \tag{15}$$

we obtain

$$\widehat{P}_{v,Q} = \left(\frac{1 + \not{k}/2m_Q}{1 - \not{k}/2m_Q}\right)^{1/2} P_{v,h}. \tag{16}$$

To determine $\widehat{P}_{v,\bar{q}}$, we can use the following useful relations

$$\left[\frac{1+\not{v}}{2}, \sqrt{\frac{1 + \not{k}/2m_Q}{1 - \not{k}/2m_Q}}\right] = -\frac{\not{k}}{2m_Q}\sqrt{\frac{1 + \not{k}/2m_Q}{1 - \not{k}/2m_Q}} \tag{17}$$

$$\sqrt{\frac{1 + \not{k}/2m_Q}{1 - \not{k}/2m_Q}} \left(\frac{1+\not{v}}{2}\right) P_{v,h} = \left(\frac{1+\not{v}}{2} + \frac{\not{k}}{2m_Q}\right) \sqrt{\frac{1 + \not{k}/2m_Q}{1 - \not{k}/2m_Q}} P_{v,h}. \tag{18}$$

Then (4) and (18) imply that $\widehat{P}_v$ relates to $P_{v,h}$ as

$$\widehat{P}_v = \sqrt{\frac{1 + \not{k}/2m_Q}{1 - \not{k}/2m_Q}} P_{v,h}. \tag{19}$$

Thus $\widehat{P}_{v,\bar{q}}$ relates to $P_{v,h}$ as

$$\widehat{P}_{v,\bar{q}} = \frac{\Lambda}{m_Q} \left(\frac{1 + \not{v}}{2} \right) \sqrt{\frac{1 + \not{k}/2m_Q}{1 - \not{k}/2m_Q}} P_{v,h}. \tag{20}$$

We can do the same procedure for the vector meson field $\widehat{P}_v^{*\mu}$ and combine all results to arrive at

$$M_v = \left(1 + \frac{\Lambda}{m_Q} \frac{1 + \not{v}}{2} \right) \sqrt{\frac{1 + \not{k}/2m_Q}{1 - \not{k}/2m_Q}} H_v. \tag{21}$$

In the above , we have arrived at the complet mass correction form for compact heavy meson effective field. It should be noted that there is no any ambiguity in this mass correction scheme. As pointed out by H.Y.Cheng, et al in[4] , there exists some ambiguities in the mass expansion via velocity reparametrization invariant method(the reason see the original paper). And such a problem has been solved partly by N.Kitazawa and T.Kurimoto by using trial and error in their derivation of mass correction lagarangian[4]. The reason about our approach to not have such a ambiguity is coming from that we directly solve the mass correction equation (4).

It needs a short comment about the meaning of the mass correction form (21). To our opinion, the key ingredient of the heavy quark mass expansion should be the field mass corrction form for the infinite mass limit field. This is because that:(1) from the renormalization group point of view , we need

an effective field theory which is continuously matching down to from the complete theory along the relevant renormalization group (**RG**) flows. In this intermediate scale region between effective and complete theories, **RG** flows are affected mostly from the heavy quark part. The mass expansion at some order should be complete under **RG** operation. This is the renormalizable problem of the heavy quark mass exapnsion. The mass expansion scheme is used to build the basis for the **RG** operation. (2) Any heavy quark involved in the green function under **RG** operation always needs one or more heavy quark fields for the calculation of its renormalization constant , both in effective and complete theories. (3)From the field theory point of view the difference between the complete and effective theory is the existence of an exotic symmetry group for effective theoy. The mass term will certainly break such symmetry. So ,the mass expansion can be seen as an transformation between effective field and original field. The conclusion is that from theoretical and practical view point the field mass correction form is most appropriate for needs.

The following we can write down the effective lagrangian which is invariant under parity transformation and charge conjugation up to $O(1/m_Q^2)$ by using (21). All the possible terms are

$$
\begin{aligned}
L &= L_{particle} + L_{antiparticle} \\
&= -\sum_v tr\{\overline{H}_v v \cdot iDH_v\} - \sum_v tr\{\overline{H}_v \frac{(i\not{D})^2}{2M} H_v\} \\
&\quad -\sum_v tr\{\overline{H}_v \frac{(i\not{D})^3}{2M^2} H_v\} + h.c.
\end{aligned}
$$

$$
\begin{aligned}
&+\Lambda\sum_v tr\{\overline{H}_v H_v\} - \Lambda\sum_v tr\{\overline{H}_v \frac{(v\cdot iD)}{2M} H_v\} + h.c.\\
&+\Lambda\sum_v tr\{\overline{H}_v \frac{(i\not{D})^2}{2M^2} H_v\} + \kappa_1\Lambda\sum_v tr\{\overline{H}_v \frac{\Lambda}{2M} H_v\}\\
&+\kappa_2\Lambda\sum_v tr\{\overline{H}_v \frac{\Lambda}{2M}\sigma_{\mu\nu} H_v\sigma^{\mu\nu}\} - \kappa_3\Lambda\sum_v tr\{\overline{H}_v \frac{\Lambda^2}{M^2} H_v\}\\
&-\kappa_4\Lambda\sum_v tr\{\overline{H}_v \frac{\Lambda^2}{2M^2}\sigma_{\mu\nu} H_v\sigma^{\mu\nu}\} + \gamma\sum_v tr\{\overline{H}_v H_v v\cdot\widehat{\alpha}_{\|}\}\\
&+\gamma\sum_v tr\{\overline{H}_v H_v v\cdot\widehat{\alpha}_{\|}\} + \gamma\sum_v tr\{\overline{H}_v \frac{iD^\mu}{2M} H_v \widehat{\alpha}_{\|\mu}\} + h.c.\\
&+\gamma\sum_v tr\{\overline{H}_v \frac{v\cdot iD}{2M} H_v v\cdot\widehat{\alpha}_{\|}\} + h.c.\\
&+\gamma_1\sum_v tr\{\overline{H}_v \frac{\Lambda}{2M} H_v v\cdot\widehat{\alpha}_{\|}\} + \gamma_2\sum_v tr\overline{H}_v \frac{\Lambda}{2M}\sigma_{\mu\nu} H\sigma^{\mu\nu} v\cdot\widehat{\alpha}_{\|}\}\\
&-\gamma_3\Lambda\sum_v tr\{\overline{H}_v \frac{5\Lambda}{2M^2} H_v v\cdot\widehat{\alpha}_{\|}\} + \gamma_4\Lambda\sum_v tr\{\overline{H}_v \frac{2iD^\mu}{2M^2} H_v\widehat{\alpha}_{\|\mu}\} + h.c.\\
&-\gamma_5\Lambda\sum_v tr\{\overline{H}_v \frac{\Lambda}{2M^2}\sigma_{\mu\nu} H\sigma^{\mu\nu} v\cdot\widehat{\alpha}_{\|}\}\\
&+\gamma\sum_v tr\{\overline{H}_v \frac{1}{2M^2}\{\Big[(iD)^2 v^\mu + (v\cdot iD) iD^\mu + iD^\mu (v\cdot iD)\Big]\} H_v\widehat{\alpha}_{\|\mu}\}\\
&+\gamma\sum_v tr\{\overline{H}_v \frac{1}{2M^2} i\epsilon^{\mu\nu\rho\sigma} iD_\mu iD_\nu H_v\gamma_5\widehat{\alpha}_{\|\rho}\gamma_\sigma\} + h.c.\\
&-\lambda\sum_v tr\{\overline{H}_v\ H_v\gamma_5\not{\alpha}_\perp\} + \lambda\sum_v tr\{\overline{H}_v \frac{v\cdot iD}{2M} H_v\gamma_5\not{\alpha}_\perp\} + h.c.\\
&-\lambda\sum_v tr\{\overline{H}_v \frac{\epsilon^{\mu\nu\rho\sigma}}{4M} iD_\rho H_v\sigma_{\mu\nu}\alpha_{\perp\sigma}\} + h.c.\\
&-\lambda_1\sum_v tr\{\overline{H}_v \frac{\Lambda}{2M} H_v\gamma_5\not{\alpha}_\perp\} - \lambda_2\sum_v tr\{\overline{H}_v \frac{\Lambda}{2M}\gamma_5\gamma_\rho H_v\alpha_\perp^\rho\}\\
&+\lambda_3\Lambda\sum_v tr\{\overline{H}_v \frac{5\Lambda}{2M^2} H_v\gamma_5\not{\alpha}_\perp\} + \lambda_4\Lambda\sum_v tr\{\overline{H}_v \frac{5\Lambda}{2M^2}\gamma_5\gamma_\rho H_v\alpha_\perp^\rho\}\\
&-\lambda_5\Lambda\sum_v tr\{\overline{H}_v \frac{2v\cdot iD}{M^2} H_v\gamma_5\not{\alpha}_\perp\} + h.c.
\end{aligned}
$$

$$
\begin{aligned}
&-\lambda_6\Lambda\sum_v tr\{\overline{H}_v\frac{5\Lambda\epsilon^{\mu\nu\rho\sigma}}{4M^2}iD_\rho H_v\sigma_{\mu\nu}\alpha_{\perp\sigma}\} + h.c.\\
&+\lambda\sum_v tr\{\overline{H}_v\frac{1}{2M^2}\{\Big[(iD)^2\gamma^\mu + (i\not{D})iD^\mu + iD^\mu(i\not{D})\Big]\}H_v\gamma_5\alpha_{\perp\mu}\}\\
&+\lambda\sum_v tr\{\overline{H}_v\frac{1}{2M^2}i\epsilon^{\mu\nu\rho\sigma}iD_\mu iD_\nu H_v\alpha_{\perp\rho}\gamma_\sigma\} + h.c.\\
&+L_{antiparticle}, \qquad (22)
\end{aligned}
$$

where $1/M = diag(1/m_c, 1/m_b)$,$1/M^2 = diag(1/m_c^2, 1/m_b^2)$. The anti-particale type lagrangian $L_{antiparticle}$ can be obtained by substituting $H_v \rightarrow H_v^-$, and $v \rightarrow -v$. For comparison , we also write down the pseudoscalar meson masses m_P and the vector meson masses m_V to be expanded as

$$
\begin{aligned}
{m_V}^2 &= {m_Q}^2\{1 + 2\frac{\Lambda}{m_Q} + \kappa_1\frac{\Lambda^2}{m_Q^2} + 6\kappa_2\frac{\Lambda^3}{m_Q^2} - 2\kappa_3\frac{\Lambda^3}{m_Q^3} - 6\kappa_4\frac{\Lambda^3}{m_Q^3}\}\\
{m_P}^2 &= {m_Q}^2\{1 + 2\frac{\Lambda}{m_Q} + \kappa_1\frac{\Lambda^2}{m_Q^2} - 2\kappa_2\frac{\Lambda^3}{m_Q^2} - 2\kappa_3\frac{\Lambda^3}{m_Q^3} + 2\kappa_4\frac{\Lambda^3}{m_Q^3}\}. \qquad (23)
\end{aligned}
$$

We adope the effective field normalization convention of [4] as $\sqrt{m_Q}$. The covariant derivative iD_μ is defined as

$$
iD_\mu H_v = i\partial_\mu H_v(x) - H_v(x)g_V V_\mu^a(x)\frac{\lambda^a}{2} \qquad (24)
$$

to include the light vector meson fields V_μ^a.

For the heavy-to-light weak lagrangian we can also obtain it in the $1/m_Q$ expansion and in the chiral expansion as

$$
L_W \;=\; i\{\frac{F}{2}tr\{J\xi^\dagger[1 + \frac{i\not{D}}{2m_Q} + \frac{(i\not{D})^2}{8m_Q^2}]H_v\}
$$

$$
\begin{aligned}
&+a_1\frac{\Lambda}{2m_Q}tr\{J\xi^\dagger H_v\}+a_2\frac{\Lambda}{2m_Q}tr\{J\xi^\dagger\gamma^\rho H_v\gamma_\rho\}\\
&+a_3\frac{\Lambda}{4m_Q^2}tr\{J\xi^\dagger v\cdot iDH_v\}+a_4\frac{\Lambda}{4m_Q^2}tr\{J\xi^\dagger v\cdot iD\gamma^\rho H_v\gamma_\rho\}\\
&+a_5\frac{\Lambda^2}{8m_Q^2}tr\{J\xi^\dagger H_v\}+a_6\frac{\Lambda^2}{8m_Q^2}tr\{J\xi^\dagger\gamma^\rho H_v\gamma_\rho\}\\
&+a_7\frac{\Lambda}{8m_Q^2}tr\{J\xi^\dagger\sigma_{\mu\nu}H_v\sigma^{\mu\nu}\}\\
&+b_1 tr\{J\xi^\dagger H_v\not{\widehat{\alpha}}_\parallel\}+b_2 tr\{J\xi^\dagger H_v v\cdot\widehat{\alpha}_\parallel\}\\
&+c_1\frac{1}{4m_Q}tr\{J\xi^\dagger iD^\mu H_v\widehat{\alpha}_{\parallel\mu}\}+c_2\frac{1}{2m_Q}tr\{J\xi^\dagger i\not{D}H_v\not{\widehat{\alpha}}_\parallel\}\\
&+c_3\frac{\Lambda}{2m_Q}tr\{J\xi^\dagger H_v\not{\widehat{\alpha}}_\parallel\}+c_4\frac{1}{8m_Q}tr\{J\xi^\dagger(i\not{D})^2H_v\not{\widehat{\alpha}}_\parallel\}\\
&+c_5\frac{\Lambda}{4m_Q^2}tr\{J\xi^\dagger v\cdot iDH_v\not{\widehat{\alpha}}_\parallel\}+c_6\frac{\Lambda^2}{8m_Q^2}tr\{J\xi^\dagger H_v\not{\widehat{\alpha}}_\parallel\}\\
&+c_7\frac{\Lambda^2}{8m_Q^2}tr\{J\xi^\dagger\sigma_{\mu\nu}H_v\sigma^{\mu\nu}\not{\widehat{\alpha}}_\parallel\}+c_8\frac{1}{2m_Q}tr\{J\xi^\dagger i\not{D}H_v v\cdot\widehat{\alpha}_\parallel\}\\
&+c_9\frac{\Lambda}{2m_Q}tr\{J\xi^\dagger H_v v\cdot\widehat{\alpha}_\parallel\}+c_{10}\frac{1}{8m_Q^2}tr\{J\xi^\dagger(i\not{D})^2H_v v\cdot\widehat{\alpha}_\parallel\}\\
&+c_{11}\frac{\Lambda}{4m_Q^2}tr\{J\xi^\dagger v\cdot iDH_v v\cdot\widehat{\alpha}_\parallel\}+c_{12}\frac{\Lambda}{4m_Q^2}tr\{J\xi^\dagger H_v v\cdot\widehat{\alpha}_\parallel\}\\
&+c_{13}\frac{\Lambda^2}{8m_Q^2}tr\{J\xi^\dagger\sigma_{\mu\nu}H_v\sigma^{\mu\nu}v\cdot\widehat{\alpha}_\parallel\}\},
\end{aligned}
\tag{25}
$$

where $J\equiv J^\mu\gamma_\mu(1-\gamma_5)$,is regarded as an external current and also invariant under parity transformation and we use the convention $J^{\mu\dagger}\equiv(Q\bar{q})$,$Q=(c,b),\bar{q}=(\bar{u},\bar{d},\bar{s})$. We can extract the decay constants of the heavy meson from the above weak lagrangian

$$
f_p=\sqrt{\frac{2}{m_Q}}\{F+\frac{\Lambda}{m_Q}(a_1+a_2)+\frac{\Lambda^2}{4m_Q^2}(a_5+a_6-4a_7)\}
\tag{26}
$$

$$f_V = \sqrt{\frac{2}{m_Q}}\{F + \frac{\Lambda}{m_Q}(a_1 - a_2) + \frac{\Lambda^2}{4m_Q^2}(a_5 - a_6 + 12a_7)\}. \tag{27}$$

We calculate the decay width of $D^{*+} \to D^o\pi^+$ as the application of our formalism. The result is

$$\begin{aligned}
&\Gamma(D^{*+}\to D^o\pi^+) \\
&\quad \approx \frac{\lambda^2 m_c^2 (E_\pi^2 - m_\pi^2)^{1/2}}{12\pi m_{D^*}^2 f_\pi^2}\{1 + E_\pi \left(\frac{1}{m_c} + \frac{5\lambda_6\Lambda^2}{m_c^2}\right) \\
&\quad + \frac{(\lambda_1 - \lambda_2)\Lambda}{\lambda m_c} - \frac{5(\lambda_3 - \lambda_4)\Lambda^2}{\lambda m_c^2}\},
\end{aligned} \tag{28}$$

where

$$E_\pi = \frac{m_{D^*}^2 - m_D^2 + m_\pi^2}{2m_{D^*}},$$

and we have dropped some double counting form factor terms.

In this paper , we have derived the mass correction of HMET by means of the projection operator method. We construct the heavy meson effective lagrangian up to $O(1/m_Q^2)$ in the $1/m_Q$ expansion and $O(p^2)$ in the chiral expansion.For the further applications we leave them to our future publications.

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